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**GENERATION OF A KM PHASE FROM SPONTANEOUS
CP BREAKING AT A HIGH ENERGY SCALE**

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ABSTRACT

We present a minimal extension of the standard model (SM) where a Kobayashi-Maskawa (KM) phase is generated from spontaneous CP breaking at a high energy scale. The only additional fields are a charge $-1/3$ vector-like quark and a singlet scalar S . In the limit $V=\sqrt{2}|<S>| \rightarrow \infty$, this model reduces to the SM with KM mechanism. For V as high as $10 - 10^2$ TeV, the neutron electric dipole moment lies in the range $10^{-26} - 10^{-28}$ ecm.

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Introduction. All present available data on CP violation [1] is consistent with having the Kobayashi-Maskawa (KM) mechanism [2] as the only source of CP violation. However, the question of the origin of the KM phase remains completely open. In gauge theories there are two distinctive ways of introducing CP violation:

- (i) Assuming that the Lagrangian violates CP through complex Yukawa couplings.
- (ii) Assuming that the Lagrangian is CP invariant, but the vacuum is not, thus leading to spontaneous CP breaking [3].

In this letter we present a minimal extension of the standard model where the KM phase is generated through spontaneous CP breaking at a high energy scale. The difficulty in achieving that has to do with the fact that in the standard $SU(2)_C \times U(1)_Y$ model with an arbitrary number of fermion families, if one imposes Natural Flavour Conservation (NFC) in the Higgs sector [4] and assumes spontaneous CP violation, then the quark-mixing matrix is real [5]. The minimal model of this class is the three Higgs doublets model suggested by Weinberg [6] in the version where CP is spontaneously violated [7].

We will consider a minimal extension of the standard model where, in the fermion sector, in addition to the three standard families, we introduce a vector-like charge $-1/3$ quark D , singlet under $SU(2)_C$, while in the Higgs sector, besides the standard doublet ϕ , we introduce a complex scalar S , singlet under $SU(2)_C \times U(1)_Y$. This minimal Higgs structure is sufficient to achieve a vacuum which spontaneously breaks CP through a phase in $\langle S \rangle$. We will show that this phase leads in general to a non-vanishing δ_{KM} in the Cabibbo-Kobayashi-Maskawa (CKM) matrix. There are flavour-changing neutral currents both in the gauge and scalar sectors but they are naturally suppressed by the small ratio $v/\sqrt{v} \ll \phi \ll S$. However, it will be shown that δ_{KM} is not suppressed by this ratio and the observed CP violation in the kaon sector can be obtained through W exchange even in the limit of arbitrarily large V . Indeed, this is a crucial feature of our model. In the limit $v \rightarrow \infty$, CP violation will be the only low energy effect left over from physics at a high energy scale. The fact that CP violation may arise exclusively from CP breaking at high energy, even in the minimal model we present, gives hope that in a grand-unified theory (GUT) one may be able to relate in a simple way [8] the observed CP violation in the kaon sector and CP violation responsible for the baryon asymmetry within the framework of the standard big bang cosmology.

The model. The quark content of the model is:

$$\begin{pmatrix} u \\ d \end{pmatrix}_L^i, u_R^i, d_R^i, D_L, \quad i = 1, 2, 3; \quad \alpha = 1, \dots, 4 \quad (1)$$

where i, α are family indices. For simplicity we have introduced only one vector-like quark, but none of the main features of the model depends on this assumption and the extension to, for example, one vector-like quark per family is straightforward. In the Higgs

sector, besides the standard model doublet, ϕ , we introduce an extra complex $SU(2)_C \times U(1)_Y$ singlet S . One may view the field S as the remnant of a GUT which breaks down to $SU(2)_C \times U(1)_Y \times SU(3)_C$. We introduce a simple Z_2 symmetry under which D_L and S are odd, while all the other fields are even. The role of the discrete symmetry is to avoid bare mass terms of the type $\bar{D}_L d_R^i$. Therefore, there will be only two mass scales in the model, namely the vacuum expectation values (vev) of ϕ and S .

We will impose CP invariance at the Lagrangian level and will show that CP can be spontaneously broken by the vacuum expectation value of the complex singlet S . The most general $SU(2)_C \times U(1)_Y \times Z_2$ invariant potential is:

$$V = V_0(\phi, S) + (\mu^2 + \lambda_1 S^* S + \lambda_2 \phi^\dagger \phi)(S^2 + S^{*2}) + \lambda_3(S^4 + S^{*4}), \quad (2)$$

where V_0 contains only terms with no phase dependence, including the standard model Higgs potential. After spontaneous symmetry breaking, one has:

$$\langle \phi^0 \rangle = v/\sqrt{2}, \quad \langle S \rangle = v e^{i\alpha}/\sqrt{2}. \quad (3)$$

The phase-dependent part of V can be written as

$$V(v, V, \alpha) = a \cos 2\alpha + b \cos 4\alpha. \quad (4)$$

It can be easily verified that for b positive and $|a| < 4b$, the absolute minimum of the potential is at:

$$\alpha = \frac{1}{2} \arccos \left(\frac{a}{4b} \right). \quad (5)$$

For this value of α , the vacuum breaks in general both T and CP invariance. It should be emphasized that spontaneous CP and T breaking occurs in this model at a high mass scale, since it is associated to the vacuum expectation value of a singlet of $SU(2)_C \times U(1)_Y$.

The gauge interactions can be written in terms of quark weak eigenstates as:

$$\mathcal{L}_{\text{gauge}} = -\frac{g}{\sqrt{2}}(\bar{u}_L^i \gamma^\mu d_L^i W_\mu^+ + h.c.) - e J_{\text{em}}^\mu A_\mu - \frac{g}{2 \cos \theta_w}(J^\mu + \bar{D}_L \gamma^\mu D_L) Z_\mu, \quad (6)$$

where

$$J^\mu = \bar{u}^i \gamma^\mu (g_V - g_A \gamma_5) u^i + \bar{d}^i \gamma^\mu (g_V - g_A \gamma_5) d^i, \quad (6')$$

is the standard model neutral current, but extended to the vector quark $d^4 \equiv D$. The crucial feature of J^μ is that it remains diagonal in flavour space when written in terms of quark mass eigenstates.

The Yukawa interactions are given by:

$$\mathcal{L}_Y = -(\bar{u} \bar{d})_L^i (g_{i\alpha} \phi d_R^\alpha + h_{ij} \phi u_R^j) - (f_\alpha S + f'_\alpha S^*) \bar{D}_L d_R^\alpha + h.c., \quad (7)$$

where $\vec{\phi} = i\sigma_2 \phi^*$, as usual. Notice that the Z_2 symmetry forbids bare mass terms for the vector quark. Upon gauge symmetry breaking, the following quark mass terms are generated:

$$\bar{u}_L^i (m_u)_{ij} u_R^j + \bar{d}_L^i (m_d)_{io} d_R^o + \bar{D}_L (M_D)_{o\alpha} d_R^\alpha + h.c., \quad (8)$$

where m_u , m_d are real 3x3 and 3x4 matrices respectively, and M_D is an arbitrary 1x4 complex matrix. We will work in a weak basis where m_u is real and diagonal. With a suitable unitary rotation of d_R^o it is possible to transform the down quark mass matrix as:

$$\mathcal{M}_d = \begin{pmatrix} m_d \\ M_D \end{pmatrix} \rightarrow \mathcal{M}'_d = \begin{pmatrix} G_{ij} & J_i \\ 0 & 0 & 0 & M \end{pmatrix}, \quad (9)$$

where M is a real number ($M = |M_D|$). The mass scale of M is set by the vev of the singlet S , but the mass scale of G and J is the weak scale v . In the decoupling limit, $M \rightarrow \infty$, the light quark mass matrix is just G and the CKM matrix is the unitary matrix that diagonalizes $GG^\dagger$. Since in the initial basis $\mathcal{M}_d$ is not the most general complex matrix one has to check the generality of $GG^\dagger$ and in particular verify that the CKM matrix can contain a δ_{KM} phase. In order to keep control of the complex phases, it is useful to make the weak basis transformation of eq. (9) through the following steps: i) with a 4x4 (real) orthogonal rotation of d_R^o one eliminates the imaginary part of $(M_D)_{i=1,2,3}$, but not of $(M_D)_4$. ii) With a 3x3 real rotation of the first three components d_R^o , M_D can be reduced to $M_D = (0 \ 0 \ M_3 \ M_4 e^{i\gamma})$ where γ is the only remaining phase in $\mathcal{M}_d$. iii) the phase shift $d_R^4 \rightarrow e^{-i\gamma} d_R^4$ eliminates the phase out of M_D but now all $(m_d)_{i4}$ are complex, with the same phase $e^{-i\gamma}$. iv) Finally, a real rotation in the $d_R^3 d_R^4$ plane of an appropriate angle ρ brings $\mathcal{M}_d$ into the desired form of eq. (9). Then,

$$\begin{aligned} G_{ik} &= m_{ik} - i p_i \tan \gamma \sin \rho \delta_{k3}, \\ J_i &= p_i / \cos \rho - m_{i3} \tan \rho - i p_i \tan \gamma \cos \rho. \end{aligned} \quad (10)$$

The angles γ , ρ and the parameters m_{ik} , p_i are all real and independent of each other. The phase γ is related to the CP violating phase of the vacuum by a relation of the type $\tan \gamma = x \tan \alpha$ where x is a ratio of scalar singlet Yukawa couplings. The matrix $\mathcal{M}'_d$ is diagonalized through the unitary transformation

$$U_L^\dagger \mathcal{M}'_d U_R = \begin{pmatrix} \bar{m} & 0 \\ 0 & M \end{pmatrix}, \quad (11)$$

where $\bar{m} = \text{diag}(m_4, m_s, m_b)$ and $\bar{M}$ is the heavy quark mass. Noting that U_L is the matrix that diagonalizes $\mathcal{M}'_d \mathcal{M}'_d{}^\dagger$, one easily derives [9] that

$$U_L = \begin{pmatrix} K & J/M \\ -J^\dagger K/M & 1 \end{pmatrix} + O\left(\frac{m^2}{M^2}\right), \quad (12)$$

where K is the unitary CKM matrix defined by: $K^\dagger G G^\dagger K = \bar{m}^2$. In addition, one obtains $M \simeq \bar{M}$. It is clear that in the limit $M \rightarrow \infty$ the standard down quarks decouple from the isosinglet D .

The CKM matrix will contain a non-vanishing δ_{KM} even in the decoupling limit. In order to check this one can use eq. (10) to verify that in general $\text{Im}\{H_{12}H_{23}H_{31}\} \neq 0$, where $H = GG^\dagger$. The crucial feature of this model is that even in the limit where the high energy scale v goes to infinity there is still mixing between all right-handed fields d_R^o . It is this mixing that generates complex phases in the mass matrix of the light quarks. However, in the decoupling limit, there is no neutral flavour violation at the weak scale because, on the one hand the left-handed isosinglet D_L decouples from the standard fields d_L^i and, on the other hand, all d_R^o have the same quantum numbers. Off decoupling limit i.e., for large but not infinite v , there are flavour violating (FV) neutral currents both through Z and Higgs exchange, but they are naturally suppressed by powers of v/Λ .

Before analysing the phenomenological implications of our model, we will derive CP invariance conditions on the quark mass matrices. This analysis will give a further insight into the mechanism of CP violation in the present model and in particular will enable us to analyse the CP properties without having to diagonalize the quark mass matrices. For the moment, we will be only interested in CP violation arising from charged and neutral gauge interactions and therefore we just consider $\mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{mass}}$, as given by eqs. (6, 8-9). One can easily derive [9] that CP invariance constrains the phases of the mass matrix $\mathcal{M}'_d$ to be of the form:

$$\arg(GG^\dagger)_{ij} = \alpha_i - \alpha_j, \quad (13)$$

where $\alpha_i = \arg J_i$. Eq. (13) shows that CP invariance completely fixes the three independent phases of $H = GG^\dagger$. In particular, $\text{Im}\{H_{12}H_{23}H_{31}\} = 0$. The total of three CP constraints equals the number of CP violating phases which in general appear in the charged and neutral currents, when one works in the quark mass eigenstates basis. So far, we have done the analysis in the weak basis of eq. (9). One may also obtain CP restrictions expressed in terms of weak basis invariants. An example is:

$$\text{Im} \text{tr}[m_u^\dagger m_d M_D^\dagger M_D m_d^\dagger m_d] = 0, \quad (14)$$

where the mass matrices are defined by eq. (8).

Phenomenological implications. We will start by writing the neutral gauge and scalar interactions in the mass eigenstates basis, giving special attention to the flavour or CP violating currents. From eqs. (6, 12) it follows that the neutral gauge interactions include the terms:

$$\mathcal{L}_Z = -\frac{g}{2\cos\theta_w} Z_\mu \left(\frac{\mu_i \mu_j^\dagger}{M^2} \bar{d}_L^i \gamma^\mu \bar{d}_L^j - \frac{\mu_i}{M} \bar{d}_L^i \gamma^\mu D_L + h.c. \right) + \dots, \quad (15)$$

where $\mu = (K^\dagger J)$, and the omitted terms are flavour conserving. Noticing that the scale of μ_i is the scale of the standard quark masses m , it follows from eq. (15) that the FV currents among light quarks are suppressed by m^2/M^2 . In order to obtain the scalar current interactions, we expand the scalar fields around the vacuum as:

$$\phi^0 = (v + h)/\sqrt{2}, \quad S = (V + s + it)e^{i\alpha}/\sqrt{2}. \quad (16)$$

Note that in the present model there are three neutral physical Higgs and no charged physical Higgs. From eq. (7) one derives the down quark Yukawa couplings in the weak basis where the mass matrix is given by eqs. (9-10):

$$\begin{aligned} \mathcal{L}_V &= -(\bar{d}_L) \begin{pmatrix} \frac{G}{v} h & \frac{f}{v} h \\ 0 & 0 \end{pmatrix} \frac{M}{V} (s + ft) \begin{pmatrix} d_R \\ s \end{pmatrix} + h.c., \\ F &= iM \sin 2\rho \frac{\sin \gamma}{\sin 2\alpha} e^{-i\gamma}, \quad f = -\tan \alpha + \frac{F}{M} \cot \rho. \end{aligned} \quad (17)$$

Only the 'transverse' component t of the scalar singlet contributes to the FV interactions because the unitary transformation that diagonalizes the down quark mass matrix also diagonalizes the couplings of the component s , 'parallel' to the vacuum. The unitary matrices U_L and U_R are defined by eq. (11) and U_L is given in leading order by eq. (12). Taking into account that U_R diagonalizes $M_d'^\dagger M_d'$ one obtains:

$$U_R = \begin{pmatrix} K_R & 0 \\ 0 & 1 \end{pmatrix} + O\left(\frac{m^2}{M^2}\right), \quad (18)$$

where K_R is defined by $K_R^\dagger G^\dagger G K_R = \bar{m}^2$. Using eqs. (12, 17-18), the relevant Yukawa couplings for the light quarks are, in leading order:

$$\begin{aligned} \mathcal{L}_V &= \left(-\frac{\bar{m}_i}{v} \delta_{ij} h + \frac{\mu_i}{v} a_j t \right) \bar{d}_L^i d_R^j + \left(-\frac{\mu_i}{v} h + \frac{\mu_i}{V} (s + ft) \right) \bar{d}_L^i D_R \\ &\quad - \left(\frac{\bar{m}_i}{v} \mu_j^* h + \frac{M}{V} a_j t \right) \bar{D}_L d_R^j + h.c., \end{aligned} \quad (19)$$

where (a_i) form a dimensionless complex vector with an absolute value $|\vec{a}|$ equal to $|F/M|$. The FV couplings between light quarks are naturally suppressed by the vev V of the singlet but not all couplings between light and heavy quarks. However, the low energy effects of these couplings are suppressed by the masses of the heavy quark and/or the heavy scalar fields.

We derive next a lower bound on the vev of the scalar singlet, V . It turns out that the most stringent bound arises from the new contributions to the neutron electric dipole moment (EDM), d_n . In lowest order, the quark d acquires an EDM, d_{dn} , from 1-loop

diagrams, as shown in fig. 1, with scalar boson exchange. Since the EDM coupling violates chirality, the result is proportional to the mass of the fermion entering in the loop. Therefore, the dominant diagrams are the ones with D , the heavy quark. In that case, denoting by M_B the mass of the virtual Higgs boson, the EDM is given in units of e by the generic expression

$$\begin{aligned} d_{dn} &= -\frac{1}{3} \frac{\text{Im}(g)}{32\pi^2 M} I\left(\frac{M^2}{M_B^2}\right), \\ I(x) &= x \frac{(x-1)(x-3) + 2\ln x}{(x-1)^3}. \end{aligned} \quad (20)$$

g stands for the product of the two Yukawa couplings to $\bar{D}_L d_R$ and $\bar{d}_L D_R$. Due to the mixing of the scalar fields, the function I has to be replaced by a sum, $V_{iB} V_{jB} I(M^2/M_B^2)$, over the mass eigenstates B , where $<ij>$ stand for a pair of fields t, h , or s , that couple to $\bar{D}_L d_R$ and $\bar{d}_L D_R$ respectively, and V is the corresponding mixing matrix. One can see in eq. (19) that the couplings $t \bar{D}_L d_R$ and $h \bar{d}_L D_R$ are not suppressed. However, the field h is essentially the standard Higgs boson H and its mixing with t, V_{tH} , is of the order of v/V because it requires the breaking of $SU(2)_C \times U(1)$. The resultant contribution is¹

$$d_{dn} \sim \pm \frac{1}{3} \frac{\text{Im}(\mu_d a_d)}{32\pi^2 V^2} [I(M^2/m_H^2) - I(M^2/M_S^2)], \quad (21)$$

where M_S stands for a typical heavy singlet mass. If one makes the reasonable assumption that m_H is smaller than M , then $I(M^2/m_H^2) \simeq 1$. Using the additive quark model i.e., $d_n = (4d_{dn} - d_{up})/3$, with $d_{up}=0$ at this order, one finally obtains, from the experimental limit $|d_n| < 10^{-25}$ cm [10], the following lower bound:

$$V > 4 \text{ TeV}, \quad (22)$$

for $|\mu_d| \sim m_s \simeq 0.2 \text{ GeV}$.

In this model there are contributions to ϵ arising from tree diagrams where a Z gauge boson or a t scalar is exchanged. It can be readily seen that the t exchange diagram is the dominant one, due to the higher suppression factor, m^2/M^2 , present in the flavour violating Z couplings. From the tree-level scalar contribution, one obtains:

$$\epsilon \simeq e^{i\pi/4} \frac{\text{Im} M_{12}}{\sqrt{2} \Delta m_K}, \quad (23)$$

$$M_{12}^S \sim (\mu_d a_s - \mu_s^* a_d)^2 \frac{f_K^2 m_K}{6M_S^2 V^2} \frac{m_K^2}{(m_s + m_d)^2} B_K. \quad (24)$$

¹ There are other contributions from the pairs $<tt>$ and $<ts>$, but they are at most of the order of magnitude of that of eq. (21).

If one takes into account the bound on V given by eq. (22), one concludes that the new contribution to ϵ is small when compared to the experimental value. Of course, the standard W exchange box diagram is also present in this model. The tree-level FV interactions also contribute to ϵ'/ϵ but, unlike the $\Delta S=2$ transition, the Z and W exchange diagrams dominate over the scalar exchange ones because both flavour conserving and flavour violating couplings of t to the light quarks are suppressed. Therefore, the new contributions to ϵ'/ϵ have the typical features present in models with vector-like quarks. According to the calculation of ref. [11], a contribution to ϵ'/ϵ of order 10^{-3} from the new diagrams would require a heavy fermion mass M as low as $3 \times 10^2 \mu$ where $\mu^2 = |\mu_4 \mu_5^*|$. In the present model, given the lower limit on V , that would require an unnaturally small ratio $M/V \sim 10^{-2}$. In addition, the present experimental limit on the $K_L \rightarrow \mu^+ \mu^-$ decay already implies $M > 3 \times 10^2 \mu/\sqrt{2}$. Thus, in our model, the only manifest deviation from the standard KM mechanism would come from a large neutron EDM.

Conclusion. We have constructed a minimal extension of the standard model, where the only two new fields are a heavy down vector quark D and a scalar singlet S . We have shown that in this minimal extension, CP can be spontaneously broken at a high mass scale, through a non-trivial phase in $< S >$. This phase in turn generates a CP violating δ_{KM} in the CKM mixing matrix of the standard three quark generations. Furthermore, this spontaneously generated δ_{KM} is not suppressed by v/V factors and indeed, for very large V (e.g. $V \sim M_{GUT} \sim 10^{15}$ GeV), δ_{KM} is the only left-over effect from spontaneous CP breaking at high energies. As a result, for very large V , the phenomenological implications of the model are identical to those of the standard model with KM mechanism. For not so large V (e.g. $V \sim 10 - 10^2$ TeV) this model distinguishes itself from the SM for a sizeable new contribution to the neutron electric dipole moment.

The connection between the observed CP violation in the kaon sector and the CP violation responsible for the generation of baryon asymmetry is an open and important question. The fact that in the simple model we have presented, one can generate CP violation at low energies from spontaneous CP breaking at a high energy, provides hope that in the ultimate GUT such a connection may be established.

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FIGURE CAPTIONS

Fig. 1 Diagram contributing to the neutron EDM.